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MIRROR WORLD AND ITS COSMOLOGICAL CONSEQUENCES

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We briefly review the concept of a parallel 'mirror' world which has the same particle physics as the observable world and couples to the latter by gravity and perhaps other very weak forces. The nucleosynthesis bounds demand that the mirror world should have a smaller temperature than the ordinary one. By this reason its evolution should substantially deviate from the standard cosmology as far as the crucial epochs like baryogenesis, nucleosynthesis etc. are concerned. In particular, we show that in the context of certain baryogenesis scenarios, the baryon asymmetry in the mirror world should be larger than in the observable one. Moreover, we show that mirror baryons could naturally constitute the dominant dark matter component of the Universe, and discuss its cosmological implications.

Keywords: Extensions of the standard model; baryogenesis; dark matter

1. Introduction

The old idea that there can exist a hidden mirror sector of particles and interactions which is the exact duplicate of our visible world [1] has attracted a significant interest over the last years. The basic concept is to have a theory given by the product $G \times G'$ of two identical gauge factors with the identical particle contents, which could naturally emerge e.g. in the context of $E_8 \times E_8'$ superstring.

In particular, one can consider a minimal symmetry $G_{\rm SM} \times G'_{\rm SM}$, where $G_{\rm SM} = SU(3) \times SU(2) \times U(1)$ stands for the standard model of observable particles: three families of quarks and leptons and the Higgs, while $G'_{\rm SM} = [SU(3) \times SU(2) \times U(1)]'$ is its mirror gauge counterpart with analogous particle content: three families of mirror quarks and leptons and the mirror Higgs. (From now on all fields and quantities of the mirror (M) sector will be marked by $'$ to distinguish from the ones belonging to the observable or ordinary (O) world.) The M-particles are singlets of $G_{\rm SM}$ and vice versa, the O-particles are singlets of $G'_{\rm SM}$. Besides the gravity, the two sectors could communicate by other means. In particular, ordinary photons could have kinetic mixing with mirror photons [2,3,4], ordinary (active) neutrinos could mix with mirror (sterile) neutrinos [5,6], or two sectors could have a common gauge symmetry of flavour [7].

A discrete symmetry $P(G \leftrightarrow G')$ interchanging corresponding fields of G and G', so called mirror parity, guarantees that both particle sectors are described by

the same Lagrangians, with all coupling constants (gauge, Yukawa, Higgs) having the same pattern, and thus their microphysics is the same.

If the mirror sector exists, then the Universe along with the ordinary photons, neutrinos, baryons, etc. should contain their mirror partners. One could naively think that due to mirror parity the ordinary and mirror particles should have the same cosmological abundances and hence the O- and M-sectors should have the same cosmological evolution. However, this would be in the immediate conflict with the Big Bang nucleosynthesis (BBN) bounds on the effective number of extra light neutrinos, since the mirror photons, electrons and neutrinos would give a contribution to the Hubble expansion rate equivalent to $\Delta N_\nu \simeq 6.14$. Therefore, in the early Universe the M-system should have a lower temperature than ordinary particles. This situation is plausible if the following conditions are satisfied:

A. After the Big Bang the two systems are born with different temperatures, namely the post-inflationary reheating temperature in the M-sector is lower than in the visible one, $T'_R < T_R$. This can be naturally achieved in certain models [8,9,10].

B. The two systems interact very weakly, so that they do not come into thermal equilibrium with each other after reheating. This condition is automatically fulfilled if the two worlds communicate only via gravity. If there are some other effective couplings between the O- and M- particles, they have to be properly suppressed.

C. Both systems expand adiabatically, without significant entropy production. If the two sectors have different reheating temperatures, during the expansion of the Universe they evolve independently and their temperatures remain different at later stages, $T' < T$, then the presence of the M-sector would not affect primordial nucleosynthesis in the ordinary world.

At present, the temperature of ordinary relic photons is $T \approx 2.75$ K, and the mass density of ordinary baryons constitutes about 5% of the critical density. Mirror photons should have smaller temperature $T' < T$, so their number density is $n'_\gamma = x^3 n_\gamma$, where $x = T'/T$. This ratio is a key parameter in our further considerations since it remains nearly invariant during the expansion of the Universe. The BBN bound on ΔN_ν implies the upper bound $x < 0.64\,\Delta N_\nu^{1/4}$. As for mirror baryons, *ad hoc* their number density n'_b can be larger than n_b, and if the ratio $\beta = n'_b/n_b$ is about 5 or so, they could constitute the dark matter of the Universe.

In this paper we study the cosmology of the mirror sector and discuss the comparative time history of the two sectors in the early Universe. We show that due to the temperature difference, in the mirror sector all key epochs as the baryogenesis, nucleosynthesis, etc. proceed at somewhat different conditions than in the observable Universe. In particular, we show that in certain baryogenesis scenarios the M-world gets a larger baryon asymmetry than the O-sector, and it is pretty plausible that $\beta > 1$.[11] This situation emerges in a particularly appealing way in the leptogenesis scenario due to the lepton number leaking from the O- to the M-sector which leads to $n'_b \geq n_b$, and can thus explain the near coincidence of visible and dark components in a rather natural way [12,13].

2. Mirror world and mirror symmetry

2.1. *Particles and couplings in the ordinary world*

Nowadays almost every particle physicist knows that particle physics is described by the Standard Model (SM) based on the gauge symmetry $G_{\rm SM} = SU(3)\times SU(2)\times U(1)$, which has a chiral fermion pattern: fermions are represented as Weyl spinors, so that the left-handed (L) quarks and leptons $\psi_L = q_L, l_L$ and right-handed (R) ones $\psi_R = q_R, l_R$ transform differently under the $SU(2)\times U(1)$ gauge factor. More precisely, the fermion content is the following:

$$l_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \sim (1,2,-1); \qquad l_R = \begin{cases} N_R \sim (1,1,0) \quad (?) \\ e_R \sim (1,1,-2) \end{cases}$$
$$q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \sim (3,2,1/3); \qquad q_R = \begin{cases} u_R \sim (3,1,4/3) \\ d_R \sim (3,1,-2/3) \end{cases}, \tag{1}$$

where the brackets explicitly indicate the $SU(3)$ and $SU(2)$ content of the multiplets and their $U(1)$ hypercharges. In addition, one prescribes a global lepton charge $L = 1$ to the leptons l_L, l_R and a baryon charge $B = 1/3$ to quarks q_L, q_R, so that baryons consisting of three quarks have $B = 1$.

The $SU(2)\times U(1)$ symmetry is spontaneously broken at the scale $v = 174$ GeV and $W^\pm, Z$ gauge bosons become massive. At the same time charged fermions get masses via the Yukawa couplings ($i,j = 1,2,3$ are the fermion generation indexes)

$$\mathcal{L}_{\rm Yuk} = Y^u_{ij}\,\overline{u_{Ri}}\,q_{Lj}\,\phi_u + Y^d_{ij}\,\overline{d_R}\,q_{Lj}\,\phi_d + Y^e_{ij}\,\overline{e_{Ri}}\,l_{Lj}\,\phi_d + \text{h.c.} \tag{2}$$

where the Higgs doublets $\phi_u = \phi \sim (1,2,1)$ and $\phi_d = \tilde{\phi} \sim (1,2,-1)$, in the minimal SM, are simply conjugated to each other: $\phi_d \sim \phi_u^*$. However, in the extensions of the SM, and in particular, in its supersymmetric version, $\phi_{u,d}$ are independent ("up" and "down") Higgs doublets with different vacuum expectation values (VEV) $\langle\phi_u\rangle = v_u$ and $\langle\phi_d\rangle = v_d$, where $v_u^2 + v_d^2 = v^2$, and their ratio is parametrized as $\tan\beta = v_u/v_d$.

Obviously, with the same rights the Standard Model could be formulated in terms of the field operators $\tilde{\psi}_R = C\gamma_0\psi_L^*$ and $\tilde{\psi}_L = C\gamma_0\psi_R^*$, where C is the charge conjugation matrix. These operators now describe antiparticles which have opposite gauge charges as well as opposite chirality than particles:

$$\tilde{l}_R = \begin{pmatrix} \tilde{\nu}_R \\ \tilde{e}_R \end{pmatrix} \sim (1,\bar{2},1); \qquad \tilde{l}_L = \begin{cases} \tilde{N}_L \sim (1,1,0) \quad (?) \\ \tilde{e}_L \sim (1,1,2) \end{cases}$$
$$\tilde{q}_R = \begin{pmatrix} \tilde{u}_R \\ \tilde{d}_R \end{pmatrix} \sim (\bar{3},\bar{2},-1/3); \qquad \tilde{q}_L = \begin{cases} \tilde{u}_L \sim (\bar{3},1,-4/3) \\ \tilde{d}_L \sim (\bar{3},1,2/3) \end{cases}. \tag{3}$$

Clearly, now antileptons $\tilde{l}_R, \tilde{l}_L$ have $L = -1$ and antiquarks $\tilde{q}_R, \tilde{q}_L$ have $B = -1/3$. The full system of particles and antiparticles can be presented in a rather symmetric way as:

$$\text{fermions}: \quad \psi_L,\ \psi_R\,; \qquad \text{antifermions}: \quad \tilde{\psi}_R,\ \tilde{\psi}_L. \tag{4}$$

However, one can simply redefine the notion of particles, namely, to call $\psi_L, \tilde{\psi}_L$ as L-particles and $\tilde{\psi}_R, \psi_R$, as R antiparticles

$$L-\text{fermions}: \quad \psi_L,\ \tilde{\psi}_L\,; \quad \tilde{R}-\text{antifermions}: \quad \tilde{\psi}_R,\ \psi_R \tag{5}$$

Hence, the standard model fields, including Higgses, can be recasted as:[a]

$$L-\text{set}: \quad (q, l, \tilde{u}, \tilde{d}, \tilde{e}, \tilde{N})_L,\ \phi_u, \phi_d\,; \quad \tilde{R}-\text{set}: \quad (\tilde{q}, \tilde{l}, u, d, e, N)_R,\ \tilde{\phi}_u, \tilde{\phi}_d \tag{6}$$

where $\tilde{\phi}_{u,d} = \phi^*_{u,d}$, and the the Yukawa Lagrangian (2) can be rewritten as

$$\mathcal{L}_{\rm Yuk} = \tilde{u}^T Y_u q \phi_u + \tilde{d}^T Y_d q \phi_d + \tilde{e}^T Y_e l \phi_d \ + \text{h.c.} \tag{7}$$

where the C-matrix as well as the family indices are omitted for simplicity.

In the absence of right-handed singlets N there are no renormalizable Yukawa couplings which could generate neutrino masses. However, once the higher order terms are allowed, the neutrinos could get Majorana masses via the $D = 5$ operators:

$$\frac{A_{ij}}{2M}(l_i\phi_2)(l_j\phi_2) \ + \ \text{h.c.} \tag{8}$$

where M is some large cutoff scale. On the other hand, introducing the N states is nothing but a natural way to generate effective operators (8) from the renormalizable couplings in the context of the seesaw mechanism. Namely, N's are gauge singlets and thus are allowed to have the Majorana mass terms $\frac{1}{2}(M_{ij}N_iN_j + M^*_{ij}\tilde{N}_i\tilde{N}_j)$. It is convenient to parametrize their mass matrix as $M_{ij} = G_{ij}M$, where M is a typical mass scale and G is a matrix of dimensionless Yukawa-like constants. On the other hand, the $\tilde{N}$ states can couple to l via Yukawa terms analogous to (7), and thus the whole set of Yukawa terms obtain the pattern:

$$l^T Y \tilde{N}\phi_u + \frac{M}{2}\tilde{N}^T G \tilde{N} + \text{h.c.} \tag{9}$$

As a result, the effective operator (8) emerges after integration out the heavy states N with $A = YG^{-1}Y^T$. This makes clear why the neutrinos masses are small – they appear in second order of the Higgs field ϕ, cutoff by large mass scale M, $m_\nu \sim v^2/M$, while the charged fermion mass terms are linear in v.

Obviously, the redefinition (6) is very convenient for the extension of the Standard Model related to supersymmetry and grand unified theories. For example, in $SU(5)$ model the L fermions fill the representations $f(\tilde{d}, l) \sim \bar{5}$ and $t(\tilde{u}, q, \tilde{e}) \sim 10$, while the R antifermions are presented by $\bar{f}(d, \tilde{l}) \sim 5$ and $\bar{t}(u, \tilde{q}, e) \sim \overline{10}$. In $SO(10)$ model all L fermions in (5) sit in one representation $L \sim 16$, while the R antifermions sit in $\tilde{R} \sim \overline{16}$.

The Standard Model crudely violates one of the possible fundamental symmetries of the Nature, parity, since its particle content and hence its Lagrangian is not

[a] In the context of $N = 1$ supersymmetry, the fermion as well as Higgs fields all become chiral superfields, and formally they can be distinguished only by matter parity Z_2 under which the fermion superfields change the sign while the Higgses remain invariant. Therefore, odd powers of what we call fermion superfields are excluded from the superpotential.

symmetric with respect to exchange of the L particles to the R ones. In particular, the gauge bosons of $SU(2)$ couple to the ψ_L fields but do not couple to ψ_R. In fact, in the limit of unbroken $SU(2) \times U(1)$ symmetry, ψ_L and ψ_R are essentially independent particles with different quantum numbers. The only reason why we call e.g. states $e_L \subset l_L$ and e_R respectively the left- and right-handed electrons, is that after the electroweak breaking down to $U(1)_{\rm em}$ these two have the same electric charges and form a massive Dirac fermion $\psi_e = e_L + e_R$.

Now what we call particles (1), the weak interactions are left-handed $(V-A)$ since only the L-states couple to the $SU(2)$ gauge bosons. In terms of antiparticles (3), the weak interactions would be right-handed $(V+A)$, since now only R states couple to the $SU(2)$ bosons. Clearly, one could always redefine the notion of particles and antiparticles, to rename particles as antiparticles and vice versa. Clearly, the natural choice for what to call particles is given by the content of matter in our Universe. Matter, at least in our galaxy and its neighbourhoods, consists of baryons q while antibaryons $\tilde{q}$ can be met only in accelerators or perhaps in cosmic rays. However, if by chance we would live in the antibaryonic island of the Universe, we would claim that our weak interactions are right-handed.

In the context of the SM or its grand unified extensions, the only good symmetry between the left and right could be the CP symmetry between L-particles and R-antiparticles. E.g., the Yukawa couplings (7) in explicit form read

$$\mathcal{L} = (\tilde{u}^T Y_u q \phi_u + \tilde{d}^T Y_d q \phi_d + \tilde{e}^T Y_e l \phi_d)_L \,+\, (u^T Y_u^* \tilde{q} \tilde{\phi}_u + d^T Y_d^* \tilde{q} \tilde{\phi}_d + e^T Y_e^* \tilde{l} \tilde{\phi}_d)_R \quad (10)$$

However, although these terms are written in an symmetric manner in terms of the L-particles and $\tilde{R}$-antiparticles, they are not invariant under $L \to \tilde{R}$ due to irremovable complex phases in the Yukawa coupling matrices. Hence, Nature does not respect the symmetry between particles and antiparticles, but rather applies the principle that "the only good discrete symmetry is a broken symmetry".

It is a philosophical question, who and how has prepared our Universe at the initial state to provide an excess of baryons over antibaryons, and therefore fixed a priority of the $V-A$ form of the weak interactions over the $V+A$ one. It is appealing to think that the baryon asymmetry itself emerges due to the CP-violating features in the particle interactions, and it is related to some fundamental physics beyond the Standard Model which is responsible for the primordial baryogenesis.

2.2. *Particles and couplings in the O- and M-worlds*

Let us assume now that there exists a mirror sector which has the same gauge group and the same particle content as the ordinary one. In the minimal version, when the O-sector is described by the gauge symmetry $G_{\rm SM} = SU(3) \times SU(2) \times U(1)$ with the observable fermions and Higgses (6), the M- sector would be given by the gauge group $G'_{\rm SM} = SU(3)' \times SU(2)' \times U(1)'$ with the analogous particle content:

$$L'\text{−set}: \; (q', l', \tilde{u}', \tilde{d}', \tilde{e}', \tilde{N}')_L, \; \phi'_u, \phi'_d; \quad \tilde{R}'\text{−set}: \; (\tilde{q}', \tilde{l}', u', d', e', N')_R, \; \tilde{\phi}'_u, \tilde{\phi}'_d \quad (11)$$

In more general view, one can consider a supersymmetric theory with a gauge symmetry $G \times G'$ based on grand unification groups as $SU(5) \times SU(5)'$, $SO(10) \times SO(10)'$, etc. The gauge factor G of the O-sector contains the vector gauge superfields V, and left chiral matter (fermion and Higgs) superfields L_a in certain representations of G, while G' stands for the M-sector with the gauge superfields V', and left chiral matter superfields $L'_a \sim r_a$ in analogous representations of G'.

The Lagrangian has a form $\mathcal{L}_{\rm gauge} + \mathcal{L}_{\rm mat}$. The matter Lagrangian is determined by the form of the superpotential which is a holomorphic function of L superfields and in general it can contain any gauge invariant combination of the latter:

$$W = (M_{ab}L_aL_b + g_{abc}L_aL_bL_c) \; + \; (M'_{ab}L'_aL'_b + g'_{abc}L'_aL'_bL'_c) + ... \tag{12}$$

where dots stand for possible higher order terms. Namely, we have

$$\mathcal{L}_{\rm mat} = \int d\theta^2 W(L) \; + \; \text{h.c.} \tag{13}$$

or, in explicit form,

$$\begin{aligned}\mathcal{L}_{\rm mat} = & \int d\theta^2 (M_{ab}L_aL_b + g_{abc}L_aL_bL_c + M'_{ab}L'_aL'_b + g'_{abc}L'_aL'_bL'_c) \\ & + \int d\overline{\theta}^2 (M^*_{ab}\tilde{R}_a\tilde{R}_b + g^*_{abc}\tilde{R}_a\tilde{R}_b\tilde{R}_c + M'^*_{ab}\tilde{R}'_a\tilde{R}'_b + g'^*_{abc}\tilde{R}'_a\tilde{R}'_b\tilde{R}'_c)\end{aligned} \tag{14}$$

where θ and $\bar{\theta}$ are the Grassmanian coordinates respectively in $(1/2, 0)$ and $(0, 1/2)$ representations of the Lorentz group, and $\tilde{R}(\tilde{R}') = L^*(L'^*)$ are the CP-conjugated right handed superfields.

One can impose a discrete symmetry between these gauge sectors in two ways:

A. Left-left symmetry D_{LL}:

$$L_a \to L'_a \quad (\tilde{R}_a \to \tilde{R}'_a), \quad V \to V'. \tag{15}$$

which for the coupling constants implies

$$M'_{ab} = M_{ab}, \qquad g'_{abc} = g_{abc} \tag{16}$$

Clearly, this is nothing but direct doubling and in this case the M-sector is an identical copy of the O-sector.

B. Left-right symmetry P_{LR}:

$$L_a \to \tilde{R}'_a \quad (\tilde{R}_a \to L'_a), \quad V \to V'. \tag{17}$$

which requires that

$$M'_{ab} = M^*_{ab}, \qquad g'_{abc} = g^*_{abc} \tag{18}$$

In this case the M-sector is a mirror copy of the ordinary sector, and this symmetry can be considered as a generalization of parity.[b]

[b]Obviously, if one imposes both D_{LL} and P_{LR}, one would get CP-invariance as a consequence.

Either type of parity implies that the two sectors have the same particle physics.[c] If the two sectors are separate and do not interact by forces other than gravity, the difference between D and P parities is rather symbolic and does not have any profound implications. However, in scenarios with some particle messengers between the two sectors this difference can be important and can have dynamical consequences.

2.3. *Couplings between O- and M-particles*

Now we discuss, what common forces could exist between the O- and M-particles, including matter fields and gauge fields.

• *Kinetic mixing term between the O- and M-photons* [2,3,4]. In the context of $G_{\rm SM} \times G'_{\rm SM}$, the general Lagrangian can contain the gauge invariant term

$$\mathcal{L} = -\chi B^{\mu\nu} B'_{\mu\nu}, \tag{19}$$

where $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$, and analogously for $B'_{\mu\nu}$, where B_μ and B'_μ are gauge fields of the abelian gauge factors $U(1)$ and $U(1)'$. Obviously, after the electroweak symmetry breaking, this term gives rise to a kinetic mixing term between the field-strength tensors of the O- and M-photons:

$$\mathcal{L} = -\varepsilon F^{\mu\nu} F'_{\mu\nu} \tag{20}$$

with $\varepsilon = \xi \cos^2\theta_W$. There is no symmetry reason for suppressing this term, and generally the constant ε could be of order 1.

On the other hand, experimental limits on the orthopositronium annihilation imply a strong upper bound on ε. This is because one has to diagonalize first the kinetic terms of the A_μ and A'_μ states and identify the physical photon as a certain linear combination of the latter. One has to notice that after the kinetic terms are brought to canonical form by diagonalization and scaling of the fields, $(A, A') \to (A_1, A_2)$, any orthonormal combination of states A_1 and A_2 becomes good to describe the physical basis. In particular, A_2 can be chosen as a "sterile" state which does not couple to O-particles but only to M-particles. Then, the orthogonal combination A_1 couples not only to O-particles, but also with M-particles with a small charge $\propto 2\varepsilon$ – in other words, mirror matter becomes "milicharged" with respect to the physical ordinary photon [2,14]. In this way, the term (20) should induce the process $e^+e^- \to e'^+e'^-$, with an amplitude just 2ε times the s-channel amplitude for $e^+e^- \to e^+e^-$. By this diagram, orthopositronium would oscillate into its mirror counterpart, which would be seen as an invisible decay mode exceeding experimental limits unless $\varepsilon < 5 \times 10^{-7}$ or so [3].

[c]The mirror parity could be spontaneously broken and the weak interaction scales $\langle\phi\rangle = v$ and $\langle\phi'\rangle = v'$ could be different, which leads to somewhat different particle physics in the mirror sector. The models with spontaneoulsy broken parity and their implications were considered in refs. [6,9,10]. In this paper we mostly concentrate on the case with exact mirror parity.

For explaining naturally the smallness of the kinetic mixing term (20) one needs to invoke the concept of grand unification. Obviously, the term (19) is forbidden if $G_{SM} \times G'_{SM}$ is embedded in GUTs like $SU(5) \times SU(5)'$ or $SO(10) \times SO(10)'$ which do not contain abelian factors. However, given that both $SU(5)$ and $SU(5)'$ symmetries are broken down to their $SU(3) \times SU(2) \times U(1)$ subgroups by the Higgs 24-plets Φ and Φ', it could emerge from the higher order effective operator

$$\mathcal{L} = -\frac{\zeta}{M^2}(G^{\mu\nu}\Phi)(G'_{\mu\nu}\Phi') \tag{21}$$

where $G_{\mu\nu}$ and $G'_{\mu\nu}$ are field-strength tensors respectively of $SU(5)$ and $SU(5)'$, and M is some cutoff scale which can be of the order of M_{Pl} or so. After substituting VEVs of Φ and Φ' the operator (20) is induced with $\varepsilon \sim \zeta(\langle\Phi\rangle/M)^2$.

In fact, the operator (21) can be effectively induced by loop-effects involving some heavy fermion or scalar fields in the mixed representations of $SU(5)$ and $SU(5)'$, with $\zeta \sim \alpha/3\pi$ being a loop-factor. Therefore, taking for the GUT scale $\langle\Phi\rangle \sim 10^{16}$ GeV and $M \sim M_{Pl}$ we see that if the kinetic mixing term (20) is induced at all, its natural range can vary from $\varepsilon \sim 10^{-10}$ up to the upper limit of 5×10^{-7}.

For $\varepsilon \sim$ few $\times 10^{-7}$ the term has striking experimental implications for positronium physics. Namely, the $e^+e^- \to e'^+e'^-$ process, would have an amplitude just 2ε times the s-channel one for $e^+e^- \to e^+e^-$, and this would lead to mixing of ordinary positronium to its mirror counterpart with significant rate, and perhaps could help in solving the troubling mismatch problems in the positronium physics [3,15,16]. However, this value is an order of magnitude above the limit $\varepsilon < 3 \times 10^{-8}$ from the BBN constraints[4]. For larger ε the reaction $e^+e^- \to e'^+e'^-$ would funnel too much energy density from the ordinary to the mirror world and would violate the BBN limit on ΔN_ν.

The search of the process $e^+e^- \to$ *invisible* could approach sensitivities down to few $\times 10^{-9}$. [17] This interesting experiment could test the proposal of ref. [18] claiming that the signal for the dark matter detection by the DAMA/NaI group [19] can be explained by elastic scattering of M-baryons with ordinary ones mediated by kinetic mixing (20), if $\varepsilon \sim 4 \times 10^{-9}$.

• *Mixing term between the O- and M-neutrinos.* In the presence of the M-sector, the $D = 5$ operator responsible for neutrino masses (8) can be immediately generalized to include an analogous terms for M-neutrinos as well as the mixed terms between the O- and M-neutrinos.

$$\frac{A_{ij}}{2M}(l_i\phi)(l_j\phi) + \frac{A'_{ij}}{2M}(l'_i\phi')(l'_j\phi') + \frac{D_{ij}}{M}(l_i\phi)(l'_j\phi') + \text{h.c.}\,, \tag{22}$$

The first operator in eq. (22), due to the ordinary Higgs vacuum VEV $\langle\phi\rangle = v \sim 100$ GeV, then induces the small Majorana masses of the ordinary (active) neutrinos. Since the mirror Higgs ϕ' also has a non-zero VEV $\langle\phi'\rangle = v'$, the second operator then provides the masses of the M-neutrinos (which in fact are sterile for the ordinary observer), and finally, the third operator induces the mixing mass terms

between the active and sterile neutrinos. The total mass matrix of neutrinos $\nu \subset l$ and $\nu' \subset l'$ reads as [6]:

$$M_\nu = \begin{pmatrix} m_\nu & m_{\nu\nu'} \\ m^T_{\nu\nu'} & m_{\nu'} \end{pmatrix} = \frac{1}{M} \begin{pmatrix} Av^2 & Dvv' \\ D^T vv' & A'v'^2 \end{pmatrix} . \tag{23}$$

Thus, this model provides a simple explanation of why sterile neutrinos could be light (on the same grounds as the active neutrinos) and could have significant mixing with the ordinary neutrinos.

The heavy neutrinos of the O-and M-sectors, $\tilde{N}$ and $\tilde{N}'$, are gauge singlets, and there is no essential difference between them. Therefore one can join states like $\tilde{N}$ and $\tilde{N}'$ in a generalized set of gauge singlet fermions N_a $(a = 1, 2, ..)$. They all could mix with each other, and, in the spirit of the seesaw mechanism, they can couple to leptons of both O- and M-sectors. The relevant Yukawa couplings have the form:

$$Y_{ia} l_i N_a \phi + Y'_{ia} l'_i N_a \phi' + \frac{1}{2} M_{ab} N_a N_b + \text{h.c.} \tag{24}$$

In this way, N play the role of messengers between ordinary and mirror particles. After integrating out these heavy states, the operators (22) are induced with $A = YG^{-1}Y^T$, $A' = Y'G^{-1}Y'^T$ and $D = YG^{-1}Y'^T$. In the next section we show that in addition the N states can mediate L and CP violating scattering processes between the O- and M-sectors which could provide a new mechanism for primordial leptogenesis.

It is convenient to present the heavy neutrino mass matrix as $M_{ab} = G_{ab}M$, M being the overall mass scale and G_{ab} some typical Yukawa constants. Without loss of generality, G_{ab} can be taken diagonal and real. Under P or D parities, in general some of the states N_a would have positive parity, while others could have a negative one.

One the other hand, the Yukawa matrices in general remain non-diagonal and complex. Then D parity would imply that $Y' = Y$, while P parity imples $Y' = Y^*$ (c.f. (16) and (18)).

• *Interaction term between the O- and M- Higgses* . In the context of $G_{\rm SM} \times G'_{\rm SM}$, the gauge symmetry allows also a quartic interaction term between the O-and M-Higgs doublets ϕ and ϕ':

$$\lambda(\phi^\dagger \phi)(\phi'^\dagger \phi') \tag{25}$$

This term is cosmologically dangerous, since it would bring the two sectors into equilibrium in the early Universe via interactions $\bar{\phi}\phi \to \bar{\phi}'\phi'$ unless λ is unnaturally small, $\lambda < 10^{-8}$.[9]

However, this term can be properly suppressed by supersymmetry. In this case standard Higgses $\phi_{u,d}$ become chiral superfields as well as their mirror partners $\phi'_{u,d}$, and so the minimal gauge invariant term between the O- and M-Higgses in the superpotential has dimension 5: $(1/M)(\phi_u\phi_d)(\phi'_u\phi'_d)$, where M is some big cutoff

mass, e.g. of the order of the GUT scale or Planck scale. Therefore, the general Higgs Lagrangian takes the form:

$$\mathcal{L} = \int d^2\theta[\mu\phi_u\phi_d + \mu\phi'_u\phi'_d + \frac{1}{M}(\phi_u\phi_d)(\phi'_u\phi'_d)] \ + \ \text{h.c.} \tag{26}$$

plus unmixed D-terms, where μ-terms are of the order of 100 GeV and M is some large cutoff scale. This Lagrangian contains a mixed quartic terms similar to (25):

$$\lambda(\phi_u^\dagger\phi_u)(\phi'_u\phi'_d) + \lambda(\phi_d^\dagger\phi_d)(\phi'_u\phi'_d) + (\phi_{u,d} \to \phi'_{u,d}) + \text{h.c.} \tag{27}$$

with the coupling constant $\lambda = \mu/M$. The same holds true for the soft supersymmetry breaking F-term and D-terms. For example, the F-term $\frac{1}{M}\int d^2\theta z(\phi_u\phi_d)(\phi'_u\phi'_d)+$ h.c., where $z = m_S\theta^2$ being the supersymmetry breaking spurion with $m_S \sim 100$ GeV, gives rise to a quartic scalar term

$$\lambda(\phi_u\phi_d)(\phi'_u\phi'_d) + \text{h.c.} \tag{28}$$

with $\lambda \sim m_S/M \ll 1$. Thus for $\mu, m_S \sim 100$ GeV, all these quartic constants are strongly suppressed, and hence are safe.

• *Mixed multiplets between the two sectors.* Until now we discussed the situation with O-particles being singlets of mirror gauge factor G' and vice versa, M-particles being singlets of ordinary gauge group G. However, in principle there could be also some fields in mixed representations of $G \times G'$. Such fields usually emerge if the two gauge factors G and G' are embedded into a bigger grand unification group $\mathcal{G}$.

For example, consider a gauge theory $SU(5) \times SU(5)'$ where the O- and M-fermions respectively are in the following left-chiral multiplets: $L \sim (\bar{5} + 10, 1)$ and $L' \sim (1, 5 + \overline{10})$. One can introduce however also the left chiral fermions in mixed representations like $F \sim (5,5)$ and $F' \sim (\bar{5},\bar{5})$, $T \sim (\overline{10},\overline{10})$ and $T' \sim (10,10)$, etc. Mixed multiplets would necessarily appear if $SU(5) \times SU(5)'$ is embbedded into $SU(10)$ group. They should have a large mass term M, e.g. of the order of $SU(10)$ breaking scale to $SU(5) \times SU(5)'$. However, in general they could couple also to the GUT Higgses $\Phi \sim (24,1)$ and $\Phi' \sim (1,24)$. Thus, their Lagrangian takes a form $MFF' + \Phi FF' + \Phi' FF' +$ h.c.

Now, under the $G_{\rm SM} \times G'_{\rm SM}$ subgroup, these multiplets split into fragments F_{ij} with different hypercharges (Y_i, Y'_j) and masses $\mathcal{M}_{ij} = M + Y_i\langle\Phi\rangle + Y'_j\langle\Phi'\rangle$. Therefore, the loops involving the fermions F_{ij} would induce a contribution to the term (19) which reads as $\chi = (\alpha/3\pi)\text{Tr}[YY'\ln(\mathcal{M}/\Lambda)]$ where Λ is an ultraviolet cutoff scale and under trace the sum over all fragments F_{ij} is understood. However, as far as these fragments emerge from the GUT multiplets, they necessarily obey that $\text{Tr}(YY') = 0$, and thus χ should be finite and cutoff independent. Thus, expanding the logarithm in terms of small parameters $\langle\Phi^{(\prime)}\rangle/M$, we finally obtain

$$\chi = \frac{\alpha\langle\Phi\rangle\langle\Phi'\rangle}{3\pi M^2}\text{Tr}[(YY')^2] \tag{29}$$

exactly what we expected from the effective operator (21). Hence, the heavy mixed multiplets in fact do not decouple and induce the O- and M-photon kinetic mixing

term proportional to the square of typical mass splittings in these multiplets ($\sim \langle\Phi\rangle^2$), analogously to the familiar situation for the photon to Z-boson mixing in the standard model.

• *Interactions via common gauge bosons.* It is pretty possible that O-and M-particles have common forces mediated by the gauge bosons of some additional symmetry group H. In other words, one can consider a theory with a gauge group $G \times G' \times H$, where O-particles are in some representations of H, $L_a \sim r_a$, and correspondingly their antiparticles are in antirepresentations, $\tilde{R}_a \sim \bar{r}_a$. As for M-particles, vice versa, we take $L'_a \sim \bar{r}_a$, and so $\tilde{R}'_a \sim r_a$. Only such a prescription of $\mathcal{G}$ pattern is compatible with the mirror parity (17). In addition, in this case H symmetry automatically becomes vector-like and so it would have no problems with axial anomalies even if the particle contents of O- and M-sectors separately are not anomaly-free with respect to H.

Let us consider the following example. The horizontal flavour symmetry $SU(3)_H$ between the quark-lepton families seems to be very promising for understanding the fermion mass and mixing pattern [21,22]. In addition, it can be useful for controlling the flavour-changing phenomena in the context of supersymmetry [7]. One can consider e.g. a GUT with $SU(5) \times SU(3)_H$ symmetry where L-fermions in (6) are triplets of $SU(3)_H$. So $SU(3)_H$ has a chiral character and it is not anomaly-free unless some extra states are introduced for the anomaly cancellation [21].

However, the concept of mirror sector makes the things easier. Consider e.g. $SU(5) \times SU(5)' \times SU(3)_H$ theory with L-fermions in (6) being triplets of $SU(3)_H$ while L'-fermions in (11) are anti-triplets. Hence, in this case the $SU(3)_H$ anomalies of the ordinary particles are cancelled by their mirror partners. Another advantage is that in a supersymmetric theory the gauge D-terms of $SU(3)_H$ are perfectly cancelled between the two sectors and hence they do not give rise to dangerous flavour-changing phenomena [7].

The immediate implication of such a theory would be the mixing of neutral O-bosons to their M-partners, mediated by horizontal gauge bosons. Namely, oscillations $\pi^0 \to \pi'^0$ or $K^0 \to K'^0$ become possible and perhaps even detectable if the horizontal symmetry breaking scale is not too high.

Another example is a common lepton number (or $B-L$) symmetry between the two sectors. Let us assume that ordinary leptons l have lepton charges $L=1$ under this symmetry while mirror ones l' have $L=-1$. Obviously, this symmetry would forbid the first two couplings in (22), $A, A' = 0$, while the third term is allowed – $D \neq 0$. Hence, 'Majorana' mass terms would be absent both for O- and M-neutrinos in (23) and so neutrinos would be Dirac particles having *naturally small* masses, with left components $\nu_L \subset l$ and right components being $\tilde{\nu}'_R \subset \tilde{l}$.

The model with common Peccei-Quinn symmetry between the O- and M-sectors was considered in [23].

3. The expansion of the Universe and thermodynamics of the O- and M-sectors

Let us assume, that after inflation ended, the O- and M-systems received different reheating temperatures, namely $T_R > T'_R$. This is certainly possible despite the fact that two sectors have identical Lagrangians, and can be naturally achieved in certain models of inflation [8,9,10].[d]

If the two systems were decoupled already after reheating, at later times t they will have different temperatures $T(t)$ and $T'(t)$, and so different energy and entropy densities:

$$\rho(t) = \frac{\pi^2}{30} g_*(T) T^4, \quad \rho'(t) = \frac{\pi^2}{30} g'_*(T') T'^4 \ , \tag{30}$$

$$s(t) = \frac{2\pi^2}{45} g_s(T) T^3, \quad s'(t) = \frac{2\pi^2}{45} g'_s(T') T'^3 \ . \tag{31}$$

The factors g_*, g_s and g'_*, g'_s accounting for the effective number of the degrees of freedom in the two systems can in general be different from each other. Let us assume that during the expansion of the Universe the two sectors evolve with separately conserved entropies. Then the ratio $x \equiv (s'/s)^{1/3}$ is time independent while the ratio of the temperatures in the two sectors is simply given by:

$$\frac{T'(t)}{T(t)} = x \cdot \left[\frac{g_s(T)}{g'_s(T')}\right]^{1/3} \ . \tag{32}$$

The Hubble expansion rate is determined by the total energy density $\bar{\rho} = \rho + \rho'$, $H = \sqrt{(8\pi/3) G_N \bar{\rho}}$. Therefore, at a given time t in a radiation dominated epoch we have

$$H(t) = \frac{1}{2t} = 1.66 \sqrt{\bar{g}_*(T)} \frac{T^2}{M_{Pl}} = 1.66 \sqrt{\bar{g}'_*(T')} \frac{T'^2}{M_{Pl}} \tag{33}$$

in terms of O- and M-temperatures $T(t)$ and $T'(t)$, where

$$\bar{g}_*(T) = g_*(T)(1 + x^4), \qquad \bar{g}'_*(T') = g'_*(T')(1 + x^{-4}). \tag{34}$$

In particular, we have $x = T'_0/T_0$, where T_0, T'_0 are the present temperatures of the O- and M- relic photons. In fact, x is the only free parameter in our model and it is constrained by the BBN bounds.

The observed abundances of light elements are in good agreement with the standard nucleosynthesis predictions. At $T \sim 1$ MeV we have $g_* = 10.75$ as it is saturated by photons γ, electrons e and three neutrino species $\nu_{e,\mu,\tau}$. The contribution of mirror particles (γ', e' and $\nu'_{e,\mu,\tau}$) would change it to $\bar{g}_* = g_*(1 + x^4)$. Deviations from $g_* = 10.75$ are usually parametrized in terms of the effective number of extra neutrino species, $\Delta g = \bar{g}_* - 10.75 = 1.75 \Delta N_\nu$. Thus we have:

$$\Delta N_\nu = 6.14 \cdot x^4 \ . \tag{35}$$

[d]For analogy, two harmonic oscillators with the same frequency (e.g. two springs with the same material and the same length) are not obliged to oscillate with the same amplitudes.

This limit very weakly depends on ΔN_ν. Namely, the conservative bound $\Delta N_\nu < 1$ implies $x < 0.64$. In view of the present observational situation, confronting the WMAP results to the BBN analysis, the bound seems to be stronger. However, e.g. $x = 0.3$ implies a completely negligible contribution $\Delta N_\nu = 0.05$.

As far as $x^4 \ll 1$, in a relativistic epoch the Hubble expansion rate (33) is dominated by the O-matter density and the presence of the M-sector practically does not affect the standard cosmology of the early ordinary Universe. However, even if the two sectors have the same microphysics, the cosmology of the early mirror world can be very different from the standard one as far as the crucial epochs like baryogenesis, nuclesosynthesis, etc. are concerned. Any of these epochs is related to an instant when the rate of the relevant particle process $\Gamma(T)$, which is generically a function of the temperature, becomes equal to the Hubble expansion rate $H(T)$. Obviously, in the M-sector these events take place earlier than in the O-sector, and as a rule, the relevant processes in the former freeze out at larger temperatures than in the latter.

In the matter domination epoch the situation becomes different. In particular, we know that ordinary baryons provide only a small fraction of the present matter density, whereas the observational data indicate the presence of dark matter with about 5 times larger density. It is interesting to question whether the missing matter density of the Universe could be due to mirror baryons? In the next section we show that this could occur in a pretty natural manner.

It can also be shown that the BBN epoch in the mirror world proceeds differently from the ordinary one, and it predicts different abundancies of primordial elements [11]. Namely, mirror helium abundance can be in the range $Y'_4 = 0.6 - 0.8$, considerably larger than the observable $Y_4 \simeq 0.24$.

4. Baryogenesis in M-sector and mirror baryons as dark matter

4.1. *Visible and dark matter in the Universe*

The present cosmological observations strongly support the main predictions of the inflationary scenario: first, the Universe is flat, with the energy density very close to the critical $\Omega = 1$, and second, primoridal density perturbations have nearly flat spectrum, with the spectral index $n_s \approx 1$. The non-relativistic matter gives only a small fraction of the present energy density, about $\Omega_m \simeq 0.27$, while the rest is attributed to the vacuum energy (cosmological term) $\Omega_\Lambda \simeq 0.73$ [24]. The fact that Ω_m and Ω_Λ are of the same order, gives rise to so called cosmological coincidence problem: why we live in an epoch when $\rho_m \sim \rho_\lambda$, if in the early Universe one had $\rho_m \gg \rho_\Lambda$ and in the late Universe one would expect $\rho_m \ll \rho_\Lambda$? The answer can be only related to an antrophic principle: the matter and vacuum energy densities scale differently with the expansion of the Universe $\rho_m \propto a^{-3}$ and $\rho_\Lambda \propto$ const., hence they have to coincide at some moment, and we are just happy to be here. Moreover, for substantially larger ρ_Λ no galaxies could be formed and thus there would not be anyone to ask this this question.

On the other hand, the matter in the Universe has two components, visible and dark: $\Omega_m = \Omega_b + \Omega_d$. The visible matter consists of baryons with $\Omega_b \simeq 0.044$ while the dark matter with $\Omega_d \simeq 0.22$ is constituted by some hypothetic particle species very weakly interacting with the observable matter. It is a tantalizing question, why the visible and dark components have so close energy densities? Clearly, the ratio

$$\beta = \frac{\rho_d}{\rho_b} \tag{36}$$

does not depend on time as far as with the expansion of the Universe both ρ_b and ρ_d scale as $\propto a^{-3}$.

In view of the standard cosmological paradigm, there is no good reason for having $\Omega_d \sim \Omega_b$, as far as the visible and dark components have different origins. The density of the visible matter is $\rho_b = M_N n_b$, where $M_N \simeq 1$ GeV is the nucleon mass, and n_b is the baryon number density of the Universe. The latter should be produced in a very early Universe by some baryogenesis mechanism, which is presumably related to some B and CP-violating physics at very high energies. The baryon number per photon $\eta = n_b/n_\gamma$ is very small. Observational data on the primordial abundances of light elements and the WMAP results on the CMBR anisotropies nicely converge to the value $\eta \approx 6 \times 10^{-10}$.

As for dark matter, it is presumably constituted by some cold relics with mass M and number density n_d, and $\rho_d = M n_d$. The most popular candidate for cold dark matter (CDM) is provided by the lightest supersymmetric particle (LSP) with $M_{\rm LSP} \sim 1$ TeV, and its number density $n_{\rm LSP}$ is fixed by its annihilation cross-section. Hence $\rho_b \sim \rho_{\rm LSP}$ requires that $n_b/n_{\rm LSP} \sim M_{\rm LSP}/M_N$ and the origin of such a conspiracy between four principally independent parameters is absolutely unclear. Once again, the value M_N is fixed by the QCD scale while $M_{\rm LSP}$ is related to the supersymmetry breaking scale, n_b is determined by B and CP violating properties of the particle theory at very high energies whereas $n_{\rm LSP}$ strongly depends on the supersymmetry breaking details. Within the parameter space of the MSSM it could vary within several orders of magnitude, and moreover, in either case it has nothing to do with the B and CP violating effects.

The situation looks even more obscure if the dark component is related e.g. to the primordial oscillations of a classic axion field, in which case the dark matter particles constituted by axions are superlight, with mass $\ll 1$ eV, but they have a condensate with enormously high number density.

In this view, the concept of mirror world could give a new twist to this problem. Once the visible matter is built up by ordinary baryons, then the mirror baryons could constitute dark matter in a natural way. They interact with mirror photons, however they are dark in terms of the ordinary photons. The mass of M-baryons is the same as the ordinary one, $M = M_N$, and so we have $\beta = n'_b/n_b$, where n'_b is the number density of M-baryons. In addition, as far as the two sectors have the same particle physics, it is natural to think that the M-baryon number density n'_b is determined by the baryogenesis mechanism which is similar to the one which fixes

the O-baryon density n_b. Thus, one could question whether the ratio $\beta = n'_b/n_b$ could be naturally order 1 or somewhat bigger.

The visible matter in the Universe consists of baryons, while the abundance of antibaryons is vanishingly small. In the early Universe, at tempreatures $T \gg 1$ GeV, the baryons and antibaryons had practically the same densities, $n_b \approx n_{\bar{b}}$ with n_b slightly exceeding $n_{\bar{b}}$, so that the baryon number density was small, $n_B = n_b - n_{\bar{b}} \ll n_b$. If there was no significant entropy production after the baryogenesis epoch, the baryon number density to entropy density ratio had to be the same as today, $B = n_B/s \approx 8 \times 10^{-11}$.[e]

One can question, who and how has prepared the initial Universe with such a small excess of baryons over antibaryons. In the Friedman Universe the initial baryon asymmetry could be arranged a priori, in terms of non-vanishing chemical potential of baryons. However, the inflationary paradigm gives another twist to this question, since inflation dilutes any preexisting baryon number of the Universe to zero. Therefore, after inflaton decay and the (re-)heating of the Universe, the baryon asymmetry has to be created by some cosmological mechanism.

There are several relatively honest baryogenesis mechanisms as are GUT baryogenesis, leptogenesis, electroweak baryogenesis, etc. (for a review, see e.g. [25]). They are all based on general principles suggested long time ago by Sakharov [26]: a non-zero baryon asymmetry can be produced in the initially baryon symmetric Universe if three conditions are fulfilled: B-violation, C- and CP-violation and departure from thermal equilibrium. In the GUT baryogenesis or leptogenesis scenarios these conditions can be satisfied in the decays of heavy particles.

At present it is not possible to say definitely which of the known mechanisms is responsible for the observed baryon asymmetry in the ordinary world. However, it is most likely that the baryon asymmetry in the mirror world is produced by the same mechanism and moreover, the properties of the B and CP violation processes are parametrically the same in both cases. But the mirror sector has a lower temperature than ordinary one, and so at epochs relevant for baryogenesis the out-of-equilibrium conditions should be easier fulfilled for the M-sector.

4.2. *Baryogenesis in the O- and M-worlds*

Let us consider the difference between the ordinary and mirror baryon asymmetries on the example of the GUT baryogenesis mechanism. It is typically based on 'slow' B- and CP-violating decays of a superheavy boson X into quarks and leptons, where slow means that at $T < M$ the Hubble parameter $H(T)$ is greater than the decay rate $\Gamma \sim \alpha M$, α being the coupling strength of X to fermions and M its mass. The other reaction rates are also of relevance: *inverse decay*: $\Gamma_I \sim \Gamma(M/T)^{3/2}\exp(-M/T)$ for $T < M_X$, and *the X boson mediated scattering*

[e] In the following we use $B = n_B/s$ which is related with the familiar $\eta = n_B/n_\gamma$ as $B \approx 0.14\eta$. However, B is more adopted for featuring the baryon asymmetry since it does not depend on time if the entropy of the Universe is conserved.

processes: $\Gamma_S \sim n_X\sigma \sim A\alpha^2T^5/M^4$, where the factor A amounts for the possible reaction channels.

The final BA depends on the temperature at which X bosons go out from equilibrium. One can introduce a parameter which measures the effectiveness of the decay at the epoch $T \sim M$: $k = (\Gamma/H)_{T=M} = 0.3\bar{g}_*^{-1/2}(\alpha M_{Pl}/M)$. For $k \ll 1$ the out-of-equilibrium condition is strongly satisfied, and per decay of one X particle one generates the baryon number proportional to the CP-violating asymmetry ε. Thus, we have $B = \varepsilon/g_*$, g_* is a number of effective degrees of freedom at $T < M$. The larger k is, the longer equilibrium is maintained and the freeze-out abundance of X boson becomes smaller. Hence, the resulting baryon number to entropy ratio becomes $B = (\varepsilon/g_*)D(k)$, where the damping factor $D(k)$ is a decreasing function of k. In particular, $D(k) = 1$ for $k \ll 1$, while for k exceeding some critical value k_c, the damping is exponential.

The presence of the mirror sector practically does not alter the ordinary baryogenesis. The effective particle number is $\bar{g}_*(T) = g_*(T)(1 + x^4)$ and thus the contribution of M-particles to the Hubble constant at $T \sim M$ is suppressed by a small factor x^4.

In the mirror sector everything should occur in a similar way, apart from the fact that now at $T' \sim M$ the Hubble constant is not dominated by the mirror species but by ordinary ones: $\bar{g}'_*(T') \simeq g'_*(T')(1 + x^{-4})$. As a consequence, we have $k' = (\Gamma/H)_{|T'=M} = kx^2$. Therefore, the damping factor for mirror baryon asymmetry can be simply obtained by replacing $k \to k' = kx^2$ in $D(k)$. In other words, the baryon number density to entropy density ratio in the M-world becomes $B' = n'_B/s' \simeq (\epsilon/g_*)D(kx^2)$. Since $D(k)$ is a decreasing function of k, then for $x < 1$ we have $D(kx^2) > D(k)$ and thus we conclude that the mirror world always gets a *larger* baryon asymmetry than the visible one, $B' > B$.[f] Namely, for $k > 1$ the baryon asymmetry in the O-sector is damped by some factor – we have $B \simeq (\varepsilon/g_*)D(k) < \varepsilon/g_*$, while if $x^2 < k^{-1}$, the damping would be irrelevant for the M-sector and hence $B' \simeq \varepsilon/g_*$.

However, this does not a priori mean that Ω'_b will be larger than Ω_b. Since the entropy densities are related as $s'/s = x^3$, for the ratio $\beta = \Omega'_b/\Omega_b$ we have:

$$\beta(x) = \frac{n'_B}{n_B} = \frac{B's'}{Bs} = \frac{x^3D(kx^2)}{D(k)} \ . \tag{37}$$

The behaviour of this ratio as a function of k for different values of the parameter x is given in the ref. [11]. Clearly, in order to have $\Omega'_b > \Omega_b$, the function $D(k)$ has to decrease faster than $k^{-3/2}$ between $k' = kx^2$ and k. Closer inspection of this function reveals that the M-baryons can be overproduced only if k is sufficiently large, so that the relevant interactions in the observable sector maintain equilibrium

[f] As it was shown in ref. [11], the relation $B' > B$ takes place also in the context of the electroweak baryogenesis scenario, where the out-of-equilibrium conditions is provided by fast phase transition and bubble nucleation.

longer than in the mirror one, and thus ordinary BA can be suppressed by an exponential Boltzmann factor while the mirror BA could be produced still in the regime $k' = kx^2 \ll 1$, when $D(k') \approx 1$.

However, the GUT baryogenesis picture has the following generic problem. In scenarios based on grand unification models like $SU(5)$, the heavy gauge or Higgs boson decays violate separately B and L, but conserve $B-L$, and so finally $B-L=0$. On the other hand, the non-perturbative sphaleron processes, which violate $B+L$ but conserve $B-L$, are effective at temperatures from about 10^{12} GeV down to 100 GeV [27]. Therefore, if $B+L$ is erased by sphaleron transitions, the final B and L both will vanish.

Hence, in a realistic scenario one actually has to produce a non-zero $B-L$ rather than just a non-zero B, a fact that strongly favours the so called *leptogenesis* scenario [28]. The seesaw mechanism for neutrino masses offers an elegant possibility of generating non-zero $B-L$ in CP-violating decays of heavy Majorana neutrinos N into leptons and Higgses. These decays violate L but obviously do not change B and so they could create a non-zero $B-L=-L_{\rm in}$. Namely, due to complex Yukawa constants, the decay rates $\Gamma(N \to l\phi)$ and $\Gamma(N \to \tilde{l}\tilde{\phi})$ can be different from each other, so that the leptons l and anti-leptons $\tilde{l}$ are produced in different amounts.

When sphalerons are in equilibrium, they violate $B+L$ and so redistribute non-zero $B-L$ between the baryon and lepton numbers of the Universe. Namely, the final values of B and $B-L$ are related as $B = a(B-L)$, where a is order 1 coefficient, namely $a \simeq 1/3$ in the SM and in its supersymmetric extension [25]. Hence, the observed baryon to entropy density ratio, $B \approx 8 \times 10^{-11}$, needs to produce $B-L \sim 2 \times 10^{-10}$.

However, the comparative analysis presented above for the GUT baryogenesis in the O- and M-worlds, is essentially true also for the leptogenesis scenario. The out-of-equilibrium decays of heavy N neutrinos of the O-sector would produce a non-zero $B-L$ which being reprocessed by sphalerons would give an observable baryon asymmetry $B = a(B-L)$. On the other hand, the same decays of heavy N' neutrinos of the M-sector would give non-zero $(B'-L')$ and thus the mirror baryon asymmetry $B' = a(B'-L')$. In order to thermally produce heavy neutrinos in both O- and M-sectors, the lightest of them should have a mass smaller than the reheating temperature T'_R in the M-sector, i.e. $M_N < T'_R, T_R$. The situation $M_N > T'_R$ would prevent thermal production of N' states, and so no $B'-L'$ would be generated in M-sector. However, one can consider also scenarios when both N and N' states are non-thermally produced in inflaton decays, but with different amounts. Then the reheating of both sectors as well as $B-L$ number generation can be related to the decays of the heavy neutrinos of both sectors and hence the situation $T'_R < T_R$ can be naturally accompanied by $B' > B$.

4.3. *Baryogenesis via Ordinary-Mirror particle exchange*

An alternative mechanism of leptogenesis based on scattering processes rather than on decay was suggested in ref. [12]. The main idea consists in the following. There exists some hidden (shadow) sector of new particles which are not in thermal equilibrium with the ordinary particle world as far as the two systems interact very weakly, e.g. if they only communicate via gravity. However, other messengers may well exist, namely, superheavy gauge singlets like right-handed neutrinos which can mediate very weak effective interactions between the ordinary and hidden leptons. Then, a net $B-L$ could emerge in the Universe as a result of CP-violating effects in the unbalanced production of hidden particles from ordinary particle collisions.

Here we consider the case when the hidden sector is a mirror one. As far as the leptogenesis is concerned, we concentrate only on the lepton sector of both O and M worlds. Therefore we consider the standard model, and among other particles species, concentrate on the lepton doublets $l_i = (\nu, e)_i$ ($i = 1, 2, 3$ is the family index) and the Higgs doublet ϕ for the O-sector, and on their mirror partners $l'_i = (\nu', e')_i$ and ϕ'. Their couplings to the heavy singlet neutrinos are given by (24).

Let us discuss now the leptogenesis mechanism in our scenario. A crucial role in our considerations is played by the reheating temperature T_R, at which the inflaton decay and entropy production of the Universe is over, and after which the Universe is dominated by a relativistic plasma of ordinary particle species. As we discussed above, we assume that after the postinflationary reheating, different temperatures are established in the two sectors: $T'_R < T_R$, i.e. the mirror sector is cooler than the visible one, or ultimately, even completely "empty".

In addition, the two particle systems should interact very weakly so that they do not come in thermal equilibrium with each other after reheating. We assume that the heavy neutrino masses are larger than T_R and thus cannot be thermally produced. As a result, the usual leptogenesis mechanism via $N \to l\phi$ decays is ineffective.

Now, the important role is played by lepton number violating scatterings mediated by the heavy neutrinos N. The "cooler" mirror world starts to be "slowly" occupied due to the entropy transfer from the ordinary sector through the $\Delta L = 1$ reactions $l_i\phi \to \bar{l}'_k\bar{\phi}'$, $\bar{l}_i\bar{\phi} \to l'_k\phi'$. In general these processes violate CP due to complex Yukawa couplings in eq. (24), and so the cross-sections with leptons and anti-leptons in the initial state are different from each other. As a result, leptons leak to the mirror sector more (or less) effectively than antileptons and a non-zero $B-L$ is produced in the Universe.

It is important to stress that this mechanism would generate the baryon asymmetry not only in the observable sector, but also in the mirror sector. In fact, the two sectors are completely similar, and have similar CP-violating properties. We have scattering processes which transform the ordinary particles into their mirror partners, and CP-violation effects in this scattering owing to the complex coupling constants. These exchange processes are active at some early epoch of the Universe,

and they are out of equilibrium. In this case, a hypothetical O observer should detect during the contact epoch that (i) matter slowly (in comparison to the Universe expansion rate) disappears from the thermal bath of our world, and, in addition, (ii) particles and antiparticles disappear with different rates, so that after the contact epoch ends up, he observes that his world is left with a non-zero baryon number even if initially it was baryon symmetric.

On the other hand, his mirror colleague, M observer, would see that (i) matter creation takes place in his world, and (ii) particles and antiparticles emerge with different rates. Therefore, after the contact epoch, he also would observe a non-zero baryon number in his world.

One would naively expect that in this case the baryon asymmetries in the O and M sectors should be literally equal, given that the CP-violating factors are the same for both sectors. However, we show that in reality, the BA in the M sector, since it is colder, can be about an order of magnitude bigger than in the O sector, as far as washing out effects are taken into account. Indeed, this effects should be more efficient for the hotter O sector while they can be negligible for the colder M sector, which could provide reasonable differences between the two worlds in case the exchange process is not too far from equilibrium. The possible marriage between dark matter and the leptobaryogenesis mechanism is certainly an attractive feature of our scheme.

The fast reactions relevant for the O-sector are the $\Delta L = 1$ one $l_i \phi \to \bar{l}'_k \bar{\phi}'$, and the $\Delta L = 2$ ones like $l\phi \to \bar{l}\bar{\phi}$, $ll \to \phi\phi$ etc. Their total rates are correspondingly

$$\Gamma_1 = \frac{Q_1 n_{\rm eq}}{8\pi M^2}; \quad Q_1 = {\rm Tr}(D^\dagger D) = {\rm Tr}[(Y'^\dagger Y')^* G^{-1} (Y^\dagger Y) G^{-1}],$$
$$\Gamma_2 = \frac{3 Q_1 n_{\rm eq}}{4\pi M^2}; \quad Q_2 = {\rm Tr}(A^\dagger A) = {\rm Tr}[(Y^\dagger Y)^* G^{-1} (Y^\dagger Y) G^{-1}], \tag{38}$$

where $n_{eq} \simeq (1.2/\pi^2) T^3$ is an equilibrium density per (bosonic) degree of freedom, and the sum is taken over all flavour and isospin indices of initial and final states. It is essential that these processes stay out of equilibrium, which means that their rates should not exceed much the Hubble parameter $H = 1.66\, g_*^{1/2} T^2/M_{Pl}$ for temperatures $T \leq T_R$, where g_* is the effective number of particle degrees of freedom, namely $g_* \simeq 100$ in the SM. In other words, the dimensionless parameters

$$k_1 = \left(\frac{\Gamma_1}{H}\right)_{T=T_R} \simeq 1.5 \times 10^{-3}\, \frac{Q_1 T_R M_{Pl}}{g_*^{1/2} M^2}$$
$$k_2 = \left(\frac{\Gamma_2}{H}\right)_{T=T_R} \simeq 9 \times 10^{-3}\, \frac{Q_2 T_R M_{Pl}}{g_*^{1/2} M^2} \tag{39}$$

should not be much larger than 1.

Let us now turn to CP-violation. In $\Delta L = 1$ processes the CP-odd lepton number asymmetry emerges from the interference between the tree-level and one-loop diagrams of fig. 1. However, CP-violation takes also place in $\Delta L = 2$ processes (see fig. 2). This is a consequence of the very existence of the mirror sector, namely, it

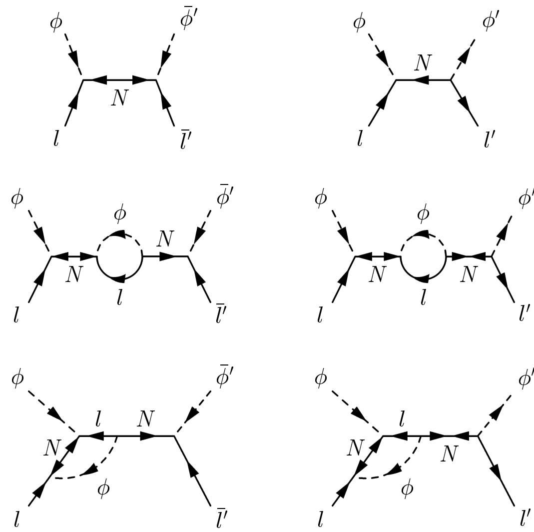

Fig. 1. Tree-level and one-loop diagrams contributing to the CP-asymmetries in $l\phi \to \bar l'\bar\phi'$ (left column) and $l\phi \to l'\phi'$ (right column).

comes from the contribution of the mirror particles to the one-loop diagrams of fig. 2. The direct calculation gives:[g]

$$
\begin{aligned}
&\sigma(l\phi \to \bar l'\bar\phi') - \sigma(\bar l\bar\phi \to l'\phi') = (-\Delta\sigma - \Delta\sigma')/2\,, \\
&\sigma(l\phi \to l'\phi') - \sigma(\bar l\bar\phi \to \bar l'\bar\phi') = (-\Delta\sigma + \Delta\sigma')/2\,, \\
&\sigma(l\phi \to \bar l\bar\phi) - \sigma(\bar l\bar\phi \to l\phi) = \Delta\sigma\,; \qquad (40)\\
&\Delta\sigma = \frac{3J\,S}{32\pi^2 M^4}\,, \qquad \Delta\sigma' = \frac{3J'\,S}{32\pi^2 M^4}\,, \qquad (41)
\end{aligned}
$$

where S is the c.m. energy square, $J = \mathrm{Im}\,\mathrm{Tr}[(Y^\dagger Y)^* G^{-1}(Y'^\dagger Y')G^{-2}(Y^\dagger Y)G^{-1}]$ is the CP-violation parameter and J' is obtained from J by exchanging Y with Y'. The contributions yielding asymmetries $\mp\Delta\sigma'$ respectively for $l\phi \to \bar l'\bar\phi'$ and $l\phi \to l'\phi'$ channels emerge from the diagrams with $l'\phi'$ inside the loops, not shown in fig. 1.

This is in perfect agreement with CPT invariance that requires that the total cross sections for particle and anti-particle scatterings are equal to each other: $\sigma(l\phi \to X) = \sigma(\bar l\bar\phi \to X)$. Indeed, taking into account that $\sigma(l\phi \to l\phi) = \sigma(\bar l\bar\phi \to \bar l\bar\phi)$ by CPT, we see that CP asymmetries in the $\Delta L = 1$ and $\Delta L = 2$ processes

[g]It is interesting to note that the tree-level amplitude for the dominant channel $l\phi \to \bar l'\bar\phi'$ goes as $1/M$ and the radiative corrections as $1/M^3$. For the channel $l\phi \to l'\phi'$ instead, both tree-level and one-loop amplitudes go as $1/M^2$. As a result, the cross section CP asymmetries are comparable for both $l\phi \to \bar l'\bar\phi'$ and $l\phi \to l'\phi'$ channels.

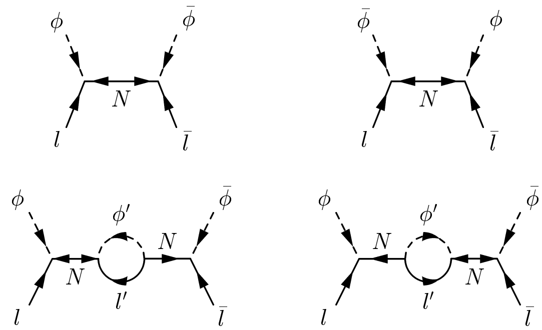

Fig. 2. Tree-level and one-loop diagrams contributing to the CP-asymmetry of $l\phi \to \bar{l}\bar{\phi}$. The external leg labels identify the initial and final state particles.

should be related as

$$\sigma(l\phi \to X') - \sigma(\bar{l}\bar{\phi} \to X') = -[\sigma(l\phi \to \bar{l}\bar{\phi}) - \sigma(\bar{l}\bar{\phi} \to l\phi)] = -\Delta\sigma \ , \tag{42}$$

where X' are the mirror sector final states, $\bar{l}'\bar{\phi}'$ and $l'\phi'$. That is, the $\Delta L = 1$ and $\Delta L = 2$ reactions have CP asymmetries with equal intensities but opposite signs.

But, as L varies in each case by a different amount, a net lepton number decrease is produced, or better, a net increase of $B - L \propto \Delta\sigma$ (recall that the lepton number L is violated by the sphaleron processes, while $B - L$ is changed solely by the above processes).

As far as we assume that the mirror sector is cooler and thus depleted of particles, the only relevant reactions are the ones with ordinary particles in the initial state. Hence, the evolution of the $B - L$ number density is determined by the CP asymmetries shown in eqs. (40) and obeys the equation

$$\frac{dn_{B-L}}{dt} + 3Hn_{B-L} + \Gamma n_{B-L} = \frac{3}{4}\Delta\sigma\, n_{\rm eq}^2 = 1.8 \times 10^{-3}\, \frac{T^8}{M^4} \ , \tag{43}$$

where $\Gamma = \Gamma_1 + \Gamma_2$ is the total rate of the $\Delta L = 1$ and $\Delta L = 2$ reactions, and for the CP asymmetric cross section $\Delta\sigma$ we take the thermal average c.m. energy square $S \simeq 17\,T^2$.

It is instructive to first solve this equation in the limit $k_{1,2} \ll 1$, when the out-of-equilibrium conditions are strongly satisfied and thus the term Γn_{B-L} can be neglected. Integrating this equation we obtain for the final $B - L$ asymmetry of the Universe, $B - L = n_{B-L}/s$, the following expression:[h]

$$(B - L)_0 \approx 2 \times 10^{-3}\, \frac{J\, M_{Pl} T_R^3}{g_*^{3/2} M^4}. \tag{44}$$

[h]Observe that the magnitude of the produced $B - L$ strongly depends on the temperature, namely, larger $B - L$ should be produced in the patches where the plasma is hotter. In the cosmological context, this would lead to a situation where, apart from the adiabatic density/temperature perturbations, there also emerge correlated isocurvature fluctuations with variable B and L which could be tested with the future data on the CMB anisotropies and large scale structure.

It is interesting to note that 3/5 of this value is accumulated at temperatures $T > T_R$ and it corresponds to the amount of $B-L$ produced when the inflaton field started to decay and the particle thermal bath was produced (Recall that the maximal temperature at the reheating period is usually larger than T_R.) In this epoch the Universe was still dominated by the inflaton oscillations and therefore it expanded as $a \propto t^{2/3}$ while the entropy of the Universe was growing as $t^{5/4}$. The other 2/5 of (44) is produced at $T < T_R$, radiation dominated era when the Universe expanded as $a \propto t^{1/2}$ with conserved entropy (neglecting the small entropy leaking from the O- to the M-sector).

This result (44) can be recasted as follows

$$(B-L)_0 \approx \frac{20Jk^2 T_R}{g_*^{1/2} Q^2 M_{Pl}} \approx 10^{-10}\, \frac{Jk^2}{Q^2} \left(\frac{T_R}{10^9 \text{ GeV}}\right) \tag{45}$$

where $Q^2 = Q_1^2 + Q_2^2$, $k = k_1 + k_2$ and we have taken again $g_* \approx 100$. This shows that for Yukawa constants spread e.g. in the range $0.1 - 1$, one can achieve $B - L = \mathcal{O}(10^{-10})$ for a reheating temperature as low as $T_R \sim 10^9$ GeV. Interestingly, this coincidence with the upper bound from the thermal gravitino production, $T_R < 4 \times 10^9$ GeV or so [30], indicates that our scenario could also work in the context of supersymmetric theories.

Let us solve now eq. (43) without assuming $\Gamma \ll H$. In this case we obtain [29]:

$$B - L = (B-L)_0 \cdot D(k) \ , \tag{46}$$

where $(B-L)_0$ is the solution of eq. (43) in the limit $\Gamma \ll H$, given by expressions (44) or (45), and the depletion factor $D(k)$ is given by

$$D(k) = \frac{3}{5}\, e^{-k} F(k) + \frac{2}{5}\, G(k) \tag{47}$$

where

$$F(k) = \frac{1}{4k^4}\left[(2k-1)^3 + 6k - 5 + 6e^{-2k}\right],$$
$$G(k) = \frac{3}{k^3}\left[2 - (k^2 + 2k + 2)e^{-k}\right]. \tag{48}$$

These two terms in $D(k)$ correspond to the integration of (43) respectively in the epochs before and after reheating ($T > T_R$ and $T < T_R$). Obviously, for $k \ll 1$ the depletion factor $D(k) \to 1$ and thus we recover the result as in (44) or (45): $B - L = (B - L)_0$. However, for large k the depletion can be reasonable, e.g. for $k = 1, 2$ we have respectively $D(k) = 0.34, 0.1$.

Now, let us discuss how the mechanism considered above produces also the baryon prime asymmetry in the mirror sector. The amount of this asymmetry will depend on the CP-violation parameter $J' = \mathrm{Im}\,\mathrm{Tr}[(Y^\dagger Y)G^{-2}(Y'^\dagger Y')Y^{-1}(Y'^\dagger Y')^* G^{-1}]$ that replaces J in $\Delta\sigma'$ of eqs. (40). The mirror P parity under the exchange $\phi \to \phi'^\dagger$, $l \to \bar{l}'$, etc., implies that the Yukawa couplings are essentially the same in both sectors, $Y' = Y^*$. Therefore, in this case

also the CP-violation parameters are the same, $J' = J$.[i] Therefore, one naively expects that $n'_{B-L} = n_{B-L}$ and the mirror baryon density should be equal to the ordinary one, $\Omega'_b = \Omega_b$.

However, now we show that if the $\Delta L = 1$ and $\Delta L = 2$ processes are not very far from equilibrium, i.e. $k \sim 1$, the mirror baryon density should be bigger than the ordinary one. Indeed, the evolution of the mirror $B-L$ number density, n'_{B-L}, obeys the equation

$$\frac{dn'_{B-L}}{dt} + 3Hn'_{B-L} + \Gamma' n'_{\rm B-L} = \frac{3}{4}\Delta\sigma'\, n^2_{\rm eq}\,, \tag{49}$$

where now $\Gamma' = (Q_1 + 6Q_2)n'_{\rm eq}/8\pi M^2$ is the total reaction rate of the $\Delta L' = 1$ and $\Delta L' = 2$ processes in the mirror sector, and $n'_{\rm eq} = (1.2/\pi^2)T'^3 = x^3 n_{\rm eq}$ is the equilibrium number density per degree of freedom in the mirror sector. Therefore $k' = \Gamma'/H = x^3 k$, and thus for the mirror sector we have $(B-L)' = (B-L)_0 D(kx^3)$, where the depletion can be irrelevant if $kx^3 \ll 1$.

Now taking into the account that in both sectors the $B-L$ densities are reprocessed into the baryon number densities by the same sphaleron processes, we have $B = a(B-L)$ and $B' = a(B-L)'$, with coefficients a equal for both sectors. Therefore, we see that the cosmological densities of the ordinary and mirror baryons should be related as

$$\beta = \frac{\Omega'_b}{\Omega_b} \approx \frac{1}{D(k)} \tag{50}$$

If $k \ll 1$, depletion factors in both sectors are $D \approx D' \approx 1$ and thus we have that the mirror and ordinary baryons have the same densities, $\Omega'_{\rm b} \approx \Omega_{\rm b}$. In this case mirror baryons are not enough to explain all dark matter and one has to invoke also some other kind of dark matter, presumably cold dark matter.

However, if $k \sim 1$, then we would have $\Omega'_{\rm b} > \Omega_{\rm b}$, and thus all dark matter of the Universe could be in the form of mirror baryons. Namely, for $k = 1.5$ we would have from eq. (50) that $\Omega'_b/\Omega_b \approx 5$, which is about the best fit relation between the ordinary and dark matter densities.

On the other hand, eq. (45) shows that $k \sim 1$ is indeed preferred for explaining the observed magnitude of the baryon asymmetry. For $k \ll 1$ the result could be too small, since $(B-L)_0 \propto k^2$ fastly goes to zero.

One could question, whether the two sectors would not equilibrate their temperatures if $k \sim 1$. As far as the mirror sector includes the gauge couplings which are the same as the standard ones, the mirror particles should be thermalized at a temperature T'. Once $k_1 \leq 1$, T' will remain smaller than the parallel temperature T of the ordinary system, and so the presence of the out-of-equilibrium hidden sector does not affect much the Big Bang nucleosynthesis epoch.

[i]It is interesting to remark that this mechanism needs the left-right parity P rather than the direct doubling one D. One can easily see that the latter requires $Y' = Y$, and so the CP-violating parameters J and J' are both vanishing.

Indeed, if the two sectors had different temperatures at reheating, then they evolve independently during the expansion of the Universe and approach the nucleosynthesis era with different temperatures. For $k_1 \leq 1$, the energy density transferred to the mirror sector will be crudely $\rho' \approx (8k_1/g_*)\rho$ [12], where $g_*(\approx 100)$ is attained to the leptogenesis epoch. Thus, translating this to the BBN limits, this corresponds to a contribution equivalent to an effective number of extra light neutrinos $\Delta N_\nu \approx k/14$.

5. Mirror baryons as dark matter

We have shown that mirror baryons could provide a significant contribution to the energy density of the Universe and thus they could constitute a relevant component of dark matter. An immediate question arises: how the mirror baryon dark matter (MBDM) behaves and what are the differences from the more familiar dark matter candidates as the cold dark matter (CDM), the hot dark matter (HDM) etc. In this section we briefly address the possible observational consequences of such a cosmological scenario.

In the most general context, the present energy density contains a relativistic (radiation) component Ω_r, a non-relativistic (matter) component Ω_m and the vacuum energy density Ω_Λ (cosmological term). According to the inflationary paradigm the Universe should be almost flat, $\Omega_0 = \Omega_m + \Omega_r + \Omega_\Lambda \approx 1$, which agrees well with the recent results on the CMBR anisotropy and large scale power spectrum.

The Hubble parameter is known to be $H_0 = 100h$ km s^{-1} Mpc^{-1} with $h \approx 0.7$, and for redshifts of cosmological relevance, $1+z = T/T_0 \gg 1$, it becomes

$$H(z) = H_0 \left[\Omega_r(1+z)^4 + \Omega_m(1+z)^3 + \Omega_\Lambda\right]^{1/2} . \tag{51}$$

In the context of our model, the relativistic fraction is represented by the ordinary and mirror photons and neutrinos, $\Omega_r h^2 = 4.2 \times 10^{-5}(1+x^4)$, and the contribution of the mirror species is negligible in view of the BBN constraint $x < 0.6$. As for the non-relativistic component, it contains the O-baryon fraction Ω_b and the M-baryon fraction $\Omega'_b = \beta\Omega_b$, while the other types of dark matter, e.g. the CDM, could also be present. Therefore, in general, $\Omega_m = \Omega_b + \Omega'_b + \Omega_{\rm cdm}$.[j]

The important moments for the structure formation are related to the matter-radiation equality (MRE) epoch and to the plasma recombination and matter-radiation decoupling (MRD) epochs.

The MRE occurs at the redshift

$$1 + z_{\rm eq} = \frac{\Omega_m}{\Omega_r} \approx 2.4 \cdot 10^4 \frac{\omega_m}{1+x^4} \tag{52}$$

[j]In the context of supersymmetry, the CDM component could exist in the form of the lightest supersymmetric particle (LSP). It is interesting to remark that the mass fractions of the ordinary and mirror LSP are related as $\Omega'_{\rm LSP} \simeq x\Omega_{\rm LSP}$. The contribution of the mirror neutrinos scales as $\Omega'_\nu = x^3\Omega_\nu$ and thus it is also irrelevant.

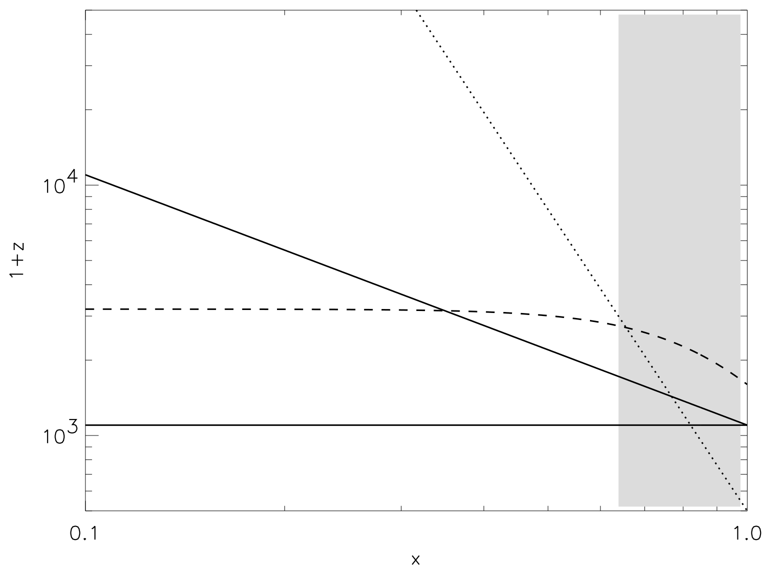

Fig. 3. The M-photon decoupling redshift $1+z'_{dec}$ as a function of x (thick solid). The horizontal thin solid line marks the ordinary photon decoupling redshift $1 + z_{\rm dec} = 1100$. We also show the matter-radiation equality redshift $1 + z_{\rm eq}$ (dash) and the mirror Jeans-horizon mass equality redshift $1 + z'_c$ (dash-dot) for the case $\omega_m = 0.135$. The shaded area $x > 0.64$ is excluded by the BBN limits.

where we denote $\omega_m = \Omega_m h^2$. Therefore, for $x \ll 1$ it is not altered by the additional relativistic component of the M-sector.

The radiation decouples from matter after almost all of electrons and protons recombine into neutral hydrogen and the free electron number density sharply diminishes, so that the photon-electron scattering rate drops below the Hubble expansion rate. In the ordinary Universe the MRD takes place in the matter domination period, at the temperature $T_{\rm dec} \simeq 0.26$ eV, which corresponds to the redshift $1 + z_{\rm dec} = T_{\rm dec}/T_0 \simeq 1100$.

The MRD temperature in the M-sector $T'_{\rm dec}$ can be calculated following the same lines as in the ordinary one [11]. Due to the fact that in either case the photon decoupling occurs when the exponential factor in Saha equations becomes very small, we have $T'_{\rm dec} \simeq T_{\rm dec}$, up to small logarithmic corrections related to B' different from B. Hence

$$1 + z'_{\rm dec} \simeq x^{-1}(1 + z_{\rm dec}) \simeq 1100\, x^{-1} \tag{53}$$

so that the MRD in the M-sector occurs earlier than in the ordinary one. Moreover, for x less than $x_{\rm eq} = 0.045\omega_m^{-1} \simeq 0.3$, the mirror photons would decouple yet during the radiation dominated period (see Fig. 3).

Let us now discuss the cosmological evolution of the MBDM. The relevant length scale for the gravitational instabilities is characterized by the mirror Jeans scale $\lambda'_J \simeq v'_s(\pi/G\rho)^{1/2}$, where $\rho(z)$ is the matter density at a given redshift z and

$v'_s(z)$ is the sound speed in the M-plasma. The latter contains more baryons and less photons than the ordinary one, $\rho'_b = \beta\rho_b$ and $\rho'_\gamma = x^4\rho_\gamma$. Let us consider for simplicity the case when dark matter of the Universe is entirely due to M-baryons, $\Omega_m \simeq \Omega'_b$. Then we have:

$$v'_s(z) \simeq \frac{c}{\sqrt{3}}\left(1+\frac{3\rho'_b}{4\rho'_\gamma}\right)^{-1/2} \approx \frac{c}{\sqrt{3}}\left[1+\frac{3}{4}\left(1+x^{-4}\right)\frac{1+z_{\rm eq}}{1+z}\right]^{-1/2}. \qquad (54)$$

Hence, for redshifts of cosmological relevance, $z \sim z_{\rm eq}$, we have $v'_s \sim 2x^2c/3 \ll c/\sqrt{3}$, quite in contrast with the ordinary world, where $v_s \approx c/\sqrt{3}$ practically until the photon decoupling, $z = 1100$.

The M-baryon Jeans mass $M'_J = \frac{\pi}{6}\rho_m\lambda'^3_J$ reaches the maximal value at $z = z'_{\rm dec} \simeq 1100/x$, $M'_J(z'_{dec}) \simeq 2.4\cdot 10^{16}\times x^6[1+(x_{\rm eq}/x)]^{-3/2}\omega_m^{-2}\ M_\odot$. Notice, however, that M'_J becomes smaller than the Hubble horizon mass $M_H = \frac{\pi}{6}\rho H^{-3}$ starting from a redshift $z_c = 3750x^{-4}\omega_m$, which is about $z_{\rm eq}$ for $x = 0.64$, but it sharply increases for smaller values of x (see Fig. 3). So, the density perturbation scales which enter the horizon at $z \sim z_{\rm eq}$ have mass larger than M'_J and thus undergo uninterrupted linear growth immediately after $t = t_{\rm eq}$. The smaller scales for which $M'_J > M_H$ instead would first oscillate. Therefore, the large scale structure formation is not delayed even if the mirror MRD epoch did not occur yet, i.e. even if $x > x_{\rm eq}$. The density fluctuations start to grow in the M-matter and the visible baryons are involved later, when after being recombined they fall into the potential whells of developed mirror structures.

Another important feature of the MBDM scenario is that the M-baryon density fluctuations should undergo strong collisional damping around the time of M-recombination. The photon diffusion from the overdense to underdense regions induce a dragging of charged particles and wash out the perturbations at scales smaller than the mirror Silk scale $\lambda'_S \simeq 3\times f(x)\omega_m^{-3/4}$ Mpc, where $f(x) = x^{5/4}$ for $x > x_{\rm eq}$, and $f(x) = (x/x_{\rm eq})^{3/2}x_{\rm eq}^{5/4}$ for $x < x_{\rm eq}$.

Thus, the density perturbation scales which can undergo the linear growth after the MRE epoch are limited by the length λ'_S. This could help in avoiding the excess of small scales (of few Mpc) in the power spectrum without tilting the spectral index. The smallest perturbations that survive the Silk damping will have the mass $M'_S \sim f^3(x)\omega_m^{-5/4}\times 10^{12}\ M_\odot$. Interestingly, for $x \sim x_{\rm eq}$ we have $M'_S \sim 10^{11}\ M_\odot$, a typical galaxy mass. To some extend, the cutoff effect is analogous to the free streaming damping in the case of warm dark matter (WDM), but there are important differences. The point is that like usual baryons, the MBDM should show acoustic oscillations whith an impact on the large scale power spectrum.

In addition, the MBDM oscillations transmitted via gravity to the ordinary baryons, could cause observable anomalies in the CMB angular power spectrum for l's larger than 200. This effect can be observed only if the M-baryon Jeans scale λ'_J is larger than the Silk scale of ordinary baryons, which sets a principal cutoff for CMB oscillations around $l \sim 1200$. As we have seen above, this would require enough large values of x, near the BBN upper bound $x \simeq 0.6$ or so.

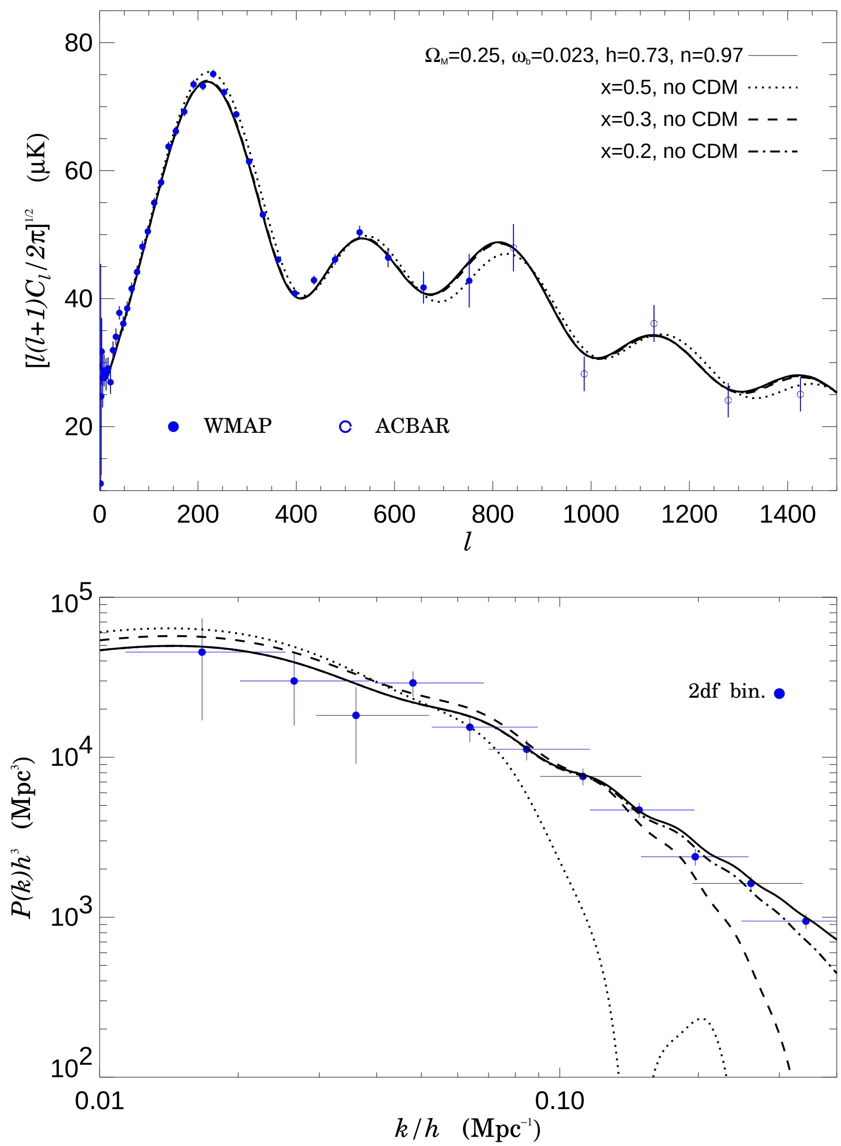

Fig. 4. The CMBR power spectrum (upper panel) and the large scale power spectrum (lower panel) for a "concordance" set of cosmological parameters. The solid curves correspond to the flat ΛCDM model, while dot, dash and dash-dot curves correspond to the situation when the CDM component is completely substituted by the MBDM for different values of x.

If the dark matter is entirely built up by mirror baryons, large values of x are excluded by the observational data. For the sake of demostration, on Fig. 4 we show the CMBR and LSS power spectra for different values of x. We see that for

$x > 0.3$ the matter power spectrum shows a strong deviation from the experimental data. This is due to Silk damping effects which suppress the small scale power too early, already for $k/h \sim 0.2$. However, the values $x < 0.3$ are compatible with the observational data.

This has a simple explanation. Clearly, for small x the M-matter recombines before the MRE moment, and thus it should rather manifest as the CDM as far as the large scale structure is concerned. However, there still can be a crucial difference at smaller scales which already went non-linear, like galaxies. Then one can question whether the MBDM distribution in halos can be different from that of the CDM? Namely, simulations show that the CDM forms triaxial halos with a density profile too clumped towards the center, and overproduce the small substructures within the halo. As for the MBDM, it constitutes a sort of collisional dark matter and thus potentially could avoide these problems, at least clearly the one related with the excess of small substructures.

As far as the MBDM constitutes a dissipative dark matter like the usual baryons, one would question how it can provide extended halos instead of being clumped into the galaxy as usual baryons do. However, one has to take into account the possibility that during the galaxy evolution the bulk of the M-baryons could fastly fragment into the stars. A difficult question to address here is related to the star formation in the M-sector, also taking into account that its temperature/density conditions and chemical contents are much different from the ordinary ones. In any case, the fast star formation would extinct the mirror gas and thus could avoide the M-baryons to form disk galaxies. The M-protogalaxy, which at a certain moment before disk formation essentially becomes a collisionless system of the mirror stars, could maintain a typical elliptical structure. In other words, we speculate on the possibility that the M-baryons form mainly elliptical galaxies.[k] Certainly, in this consideration also the galaxy merging process should be taken into account. As for the O-matter, within the dark M-matter halo it should typically show up as an observable elliptic or spiral galaxy, but some anomalous cases can also be possible, like certain types of irregular galaxies or even dark galaxies dominantly made of M-baryons.

Another tempting issue is whether the M-matter itself could help in producing big central black holes, with masses up to $\sim 10^9\ M_\odot$, which are thought to be the main engines of active galactic nuclei.

Another possibility can also be considered when dark matter in galaxies and clusters contain mixed CDM and MBDM components, $\Omega_d = \Omega'_b + \Omega_{cdm}$. e.g. one can exploit the case when mirror baryons constitute the same fraction of matter as the ordinary ones, $\Omega'_b = \Omega_b$, a situation which emerges naturally in the leptogenesis mechanism of sect. 4.3 in the case of small k.

In this case the most interesting and falsificable predictions are related to the

[k]For a comparison, in the ordinary world the number of spiral and elliptic galaxies are comparable. Remarkably, the latter contain old stars, very little dust and show low activity of star formation.

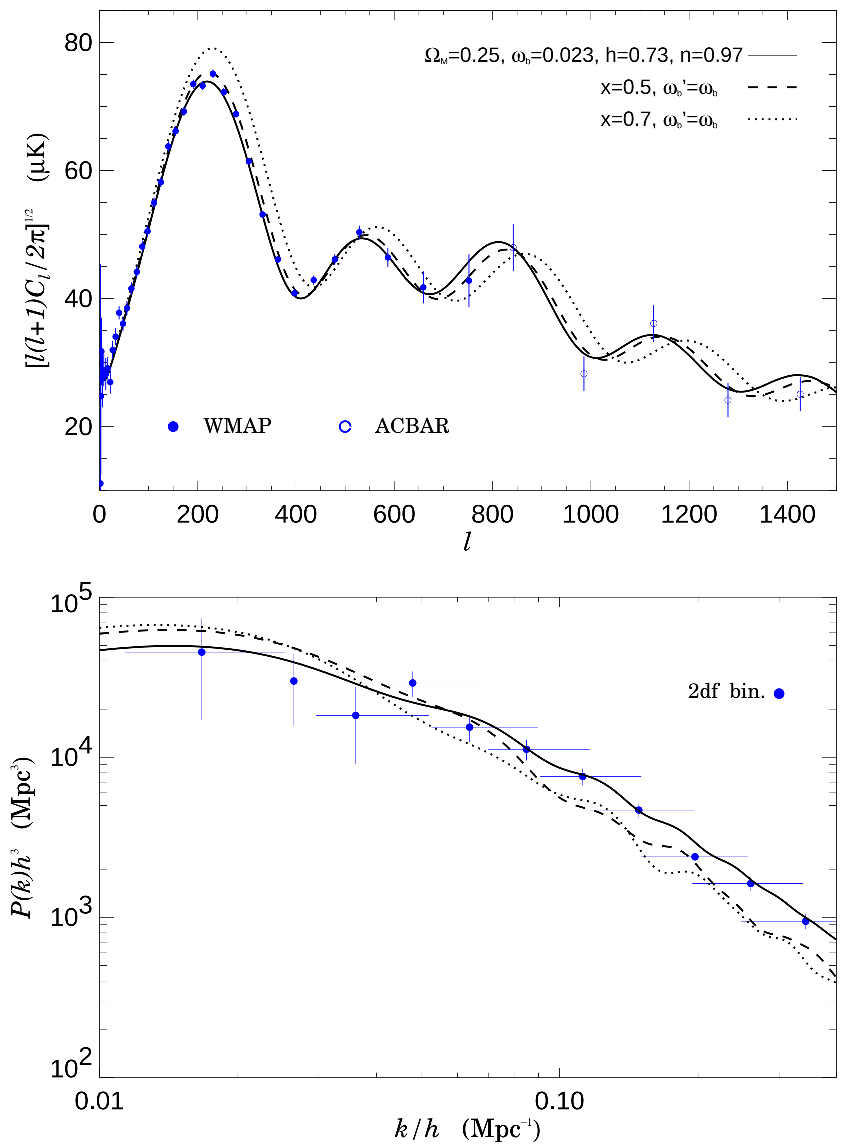

Fig. 5. The same as on Fig. 4, however for the mixed CDM+MBDM scenario for large values of x. The ordinary and mirror baryon densities are taken equal, $\Omega'_b = \Omega_B$, and the rest of matter density is attained to the SDM component.

large x regime. On Fig. 5 we show the results for the CMBR and LSS power spectra. We see that too large values of x are excluded by the CMBR anisotropies, but e.g. $x \leq 0.5$ can still be compatible with the data.

The detailed analysis of this effect will be given elsewhere [31]. In our opinion, in case of large x the effects on the CMBR and LSS can provide direct tests for the

MBDM and can be falsified by the next observations with higher sensitivity.

In the galactic halo (provided that it is an elliptical mirror galaxy) the mirror stars should be observed as Machos in gravitational microlensing [9,32]. Leaving aside the difficult question of the initial stellar mass function, one can remark that once the mirror stars could be very old and evolve faster than the ordinary ones, it is suggestive to think that most of the massive ones, with mass above the Chandrasekhar limit $M_{\rm Ch} \simeq 1.5\ M_\odot$, have already ended up as supernovae, so that only the lighter ones remain as the microlensing objects. The recent data indicate the average mass of Machos around $M \simeq 0.5\ M_\odot$, which is difficult to explain in terms of the brown dwarves with masses below the hydrogen ignition limit $M < 0.1 M_\odot$ or other baryonic objects [33]. Perhaps, this is the observational evidence of mirror matter?

It is also possible that in the galactic halo some fraction of mirror stars exists in the form of compact substructures like globular or open clusters. In this case, for a significant statistics, one could observe interesting time and angular correlations between the microlensing events.

The explosions of mirror supernovae in our galaxy cannot be directly seen by an ordinary observer. However, it should be observed in terms of gravitational waves. In addition, if the M- and O-neutrinos are mixed [5,6], it can lead to an observable neutrino signal, and could be also accompanied by a weak gamma ray burst [34].

6. Conclusions and outlook

We have discussed cosmological implications of the parallel mirror world with the same microphysics as the ordinary one, but having smaller temperature, $T' < T$, with the limit on $x = T'/T < 0.6$ set by the BBN constraints. Therefore, the M-sector contains less relativistic matter (photons and neutrinos) than the O-sector, $\Omega'_r \ll \Omega_r$. On the other hand, in the context of certain baryogenesis scenarios, the condition $T' < T$ yields that the mirror sector should produce a larger baryon asymmetry than the observable one, $B' > B$. So, in the relativistic expansion epoch the cosmological energy density is dominated by the ordinary component, while the mirror one gives a negligible contribution. However, for the non-relativistic epoch the complementary situation can occur when the mirror baryon density is bigger than the ordinary one, $\Omega'_b > \Omega_b$. Hence, the MBDM can contribute as dark matter along with the CDM or even entirely constitute it.

Unfortunately, we cannot exchange the information with the mirror physicists and combine our observations. (After all, since the two worlds have the same microphysics, life should be possible also in the mirror sector.) However, there can be many possibilities to disentangle the cosmological scenario of two parallel worlds with the future high precision data concerning the large scale structure, CMB anisotropy, structure of the galaxy halos, gravitational microlensing, oscillation of neutrinos or other neutral particles into their mirror partners, etc.

Let us conclude with two quotes of a renowned theorist. In 1986 Glashow found

a contradiction between the estimates of the GUT scale induced kinetic mixing term (20) and the positronium limits $\varepsilon \leq 4 \times 10^{-7}$ and concluded that [3]: "*Since these are in evident conflict, the notion of a mirror universe with induced electromagnetic couplings of plausible (or otherwise detectable) magnitudes is eliminated. The unity of physics is again demonstrated when the old positronium workhorse can be recalled to exclude an otherwise tenable hypothesis*".

The situation got another twist within one year, after the value $\varepsilon \approx 4 \times 10^{-7}$ appeared to be just fine for tackling the mismatch problem of the orthopositronium lifetime. However, in 1987 Glashow has fixed that this value was in conflict with the BBN limit $\varepsilon < 3 \times 10^{-8}$ and concluded the following [4]: "*We see immediately that this limit on ϵ excludes mirror matter as an explanation of the positronium lifetime ... We also note that the expected range for ϵ $(10^{-3} - 10^{-8})$ is also clearly excluded. This suggests that the mirror universe, if it exists at all, couples only gravitationally to our own. If the temperature of the mirror universe is much lower than our own, then no nucleosynthesis limit can be placed on the mirror universe at all. Then it is also likely that the mirror universe would have a smaller baryon number as well, and hence would be virtually empty. This makes a hypothetical mirror universe undetectable at energies below the Planck energy. Such a mirror universe can have no influence on the Earth and therefore would be useless and therefore does not exist*".

The main purpose of this paper was to object to this statement. The mirror Universe, if it exists at all, would be useful and can have an influence if not directly on the Earth, but on the formation of galaxies ... and moreover, the very existence of matter, both of visible and dark components, can be a consequence of baryogenesis via entropy exchange between the two worlds. The fact that the temperature of the mirror Universe is much lower than the one in our own, does not imply that it would have a smaller baryon number as well and hence would be virtually empty, but it is likely rather the opposite, mirror matter could have larger baryon number and being more matter-rich, it can provide a plausible candidate for dark matter in the form of mirror baryons. Currently it seems to be the only concept which could naturally explain the coincidence between the visible and dark matter densities of the Universe. In this view, future experiments for direct detection of mirror matter are extremely interesting.

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